

The LHC diphoton resonance from gauge symmetry

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Motivated by what is possibly the first sign of new physics seen at the LHC, the diphoton excess at 750 GeV in ATLAS and CMS, we present a model that provides naturally the necessary ingredients to explain the resonance. The simplest phenomenological explanation for the diphoton excess requires a new scalar state, $X(750)$, as well as additional vector-like (VL) fermions introduced in an ad-hoc way in order to enhance its decays into a pair of photons and/or increase its production cross-section. We show that the necessary VL quarks and their couplings can emerge naturally from a complete framework based on the $SU(3)_L \otimes U(1)_X$ gauge symmetry.

I. INTRODUCTION AND MOTIVATIONS

The great expectations to find New Physics (NP) at the LHC may have materialized with the observation of a diphoton excess at $\sim 4\sigma$ at ~ 750 GeV by the ATLAS and CMS collaborations [1, 2]. This signal, if confirmed by further data, would be the first clear and direct LHC indication of physics beyond the Standard Model (SM), a framework that was anyway expected to be incomplete due to its failure to account for non-zero neutrino masses, dark matter, and the matter-antimatter asymmetry of the Universe.

Given the low statistical significance of this hint (locally 3.6σ and 2.6σ in ATLAS and CMS, respectively), it is definitely too soon to claim the end of the SM. Nevertheless, it is tempting to speculate about the possible origin of such an excess. In fact, the announcement of this hint has triggered an intense activity translated into many recent papers analyzing the diphoton excess and proposing various physical explanations as to its origin [3–63]. Generally speaking, the simplest new physics interpretation of the diphoton excess is through the radiative decay of a new spin-0 state produced resonantly at the LHC.

In what follows we will assume that indeed NP is at work here and the resonance, to which we refer as $X(750)$, is genuine. In contrast with most explanations to the diphoton excess proposed so far, which introduce new ad-hoc states to enhance the diphoton rate or increase the $X(750)$ production cross-section, we will contemplate the possibility that this particle, as well as

the necessary ingredients to get the required diphoton signal, are the result of some gauge extension of the SM.

Perhaps, the simplest phenomenological extension of the SM that can account for $X(750)$ is the addition of a real scalar singlet, that we denote as X , and a vector-like (VL) quark, Q . The combination of these two elements allows us to write a phenomenological Lagrangian,

$$-\mathcal{L}_{\text{pheno}} = \frac{1}{2}M_X X^2 + M_Q \bar{Q}Q + \lambda X \bar{Q}Q, \quad (1)$$

which effectively generates the interactions with gluons and photons, $c_s X G_{\mu\nu} G^{\mu\nu} + c_e X F_{\mu\nu} F^{\mu\nu}$, with $c_a \propto \frac{\lambda \alpha_a^2}{M_Q}$ for $a = s, e$ and $\alpha_{s(e)}$ is the strong (electromagnetic) coupling strength.

It is our goal here to generate such effective interactions from a gauge extension of the SM. A simple embedding of $SU(2)_L \otimes U(1)_Y$ into a larger group is provided by $SU(3)_L \otimes U(1)_X$. The group structure forces the introduction of new colored fermions to complete the $SU(3)_L$ multiplets. These new quarks are $SU(2)_L \otimes U(1)_Y$ singlets after the breaking of $SU(3)_L \otimes U(1)_X$, and offer the attractive possibility to account for eq. (1) from the gauge symmetry. Indeed, if we take for instance a singlet right-handed quark field Q_R , a “weak” quark multiplet, and a scalar in the fundamental representations of $SU(3)_L$:

$$Q_L = \begin{pmatrix} q \\ Q \end{pmatrix}_L, \quad \Phi = \begin{pmatrix} H \\ X \end{pmatrix}, \quad (2)$$

which after the breaking $SU(3)_L \otimes U(1)_X \rightarrow SU(2)_L \otimes U(1)_Y$ gives the quarks and Higgs $SU(2)_L$ doublets, q_L and H , as well as the iso-singlets Q and X . We see that the gauge-invariant coupling $\Phi \bar{Q}_L Q_R$ automatically generates at low energies the coupling $X \bar{Q}Q$ as required in eq. (1).

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In addition to this attractive feature, models based on $SU(3)_c \otimes SU(3)_L \otimes U(1)_\mathcal{X}$ gauge symmetry (3-3-1 for short) [64–67] constitute a minimal extension of the SM that could explain the number of generations, and provide mechanisms to generate small neutrino masses either radiatively [68, 69] or at tree-level with new testable flavor predictions [70] and gauge bosons physics lying at the TeV scale. This can also be related to gauge coupling unification [71] and interesting D-brane constructions [72].

The paper is organized as follows. In the next Section we present a complete $SU(3)_L \otimes U(1)_\mathcal{X}$ model with all the ingredients to explain the diphoton hint observed by ATLAS and CMS. In Sec. III we derive the low energy Lagrangian after $SU(3)_L \otimes U(1)_\mathcal{X}$ breaking, whereas in Sec. IV we show how this setup naturally accommodates the $X(750)$ state in a straightforward and natural way, thus providing a complete framework for the diphoton anomaly. Finally, we will conclude with a discussion.

II. THE MODEL

We consider a variant of the models in [68, 70]. The model is based on the $SU(3)_c \otimes SU(3)_L \otimes U(1)_\mathcal{X}$ gauge symmetry and contains three generations of lepton triplets (ψ_L), two generations of quark triplets ($Q_L^{1,2}$), one generation of quark anti-triplet (Q_L^3), plus their iso-singlet right-handed partners. The gauge symmetry breaking is achieved through four scalar anti-triplets ($\Phi_{1,2,3}$ and Φ_X). In our conventions, the electric charge generator is defined as $Q = T_3 + \frac{1}{\sqrt{3}}T_8 + \mathcal{X}$ where $T_{3,8}$ are the diagonal generators of $SU(3)_L$.¹ The particle content of the model is summarized in table I.

The fermions representations of the model can be decomposed as:

$$\psi_L = \begin{pmatrix} \ell^- \\ -\nu \\ N^c \end{pmatrix}_L^{e,\mu,\tau}, \quad Q_L^1 = \begin{pmatrix} u \\ d \\ D \end{pmatrix}_L, \quad Q_L^2 = \begin{pmatrix} c \\ s \\ S \end{pmatrix}_L, \quad Q_L^3 = \begin{pmatrix} b \\ -t \\ T \end{pmatrix}_L \quad (3)$$

The notation used for the extra quarks that constitute the third components of the $SU(3)_L$ triplets $Q_L^{1,2,3}$ is motivated by the fact that their electric charges are $-1/3$ and $2/3$ for D/S and T , respectively. The scalar multi-

	ψ_L	l_R	$Q_L^{1,2}$	Q_L^3	q_R^u	T_R	q_R^d	D_R, S_R	Φ_1	Φ_2	Φ_3	Φ_X
$SU(3)_C$	1	1	3	3	3	3	3	3	1	1	1	1
$SU(3)_L$	3*	1	3	3*	1	1	1	1	3*	3*	3*	3*
$U(1)_\mathcal{X}$	$-\frac{1}{3}$	-1	0	$+\frac{1}{3}$	$+\frac{2}{3}$	$+\frac{2}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$	$+\frac{2}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$
$U(1)_\mathcal{L}$	$-\frac{1}{3}$	-1	$-\frac{2}{3}$	$+\frac{2}{3}$	0	0	0	0	$+\frac{2}{3}$	$-\frac{4}{3}$	$+\frac{2}{3}$	$-\frac{4}{3}$
\mathbf{Z}_2	+	+	+	-	+	-	-	+	+	+	-	+

TABLE I. Particle content of the model. Here $q_R^u \equiv \{u_R, c_R, t_R\}$, $q_R^d \equiv \{d_R, s_R, b_R\}$. The global symmetry $U(1)_\mathcal{L}$ allows for a consistent definition of lepton number via the relation $L = \frac{4}{\sqrt{3}}T_8 + \mathcal{L}$ and the \mathbf{Z}_2 parity simplifies the expressions of quark masses. See Ref. [68] for further details.

plets can be written as:

$$\Phi_1 = \begin{pmatrix} \phi_1^+ \\ -\phi_1^+ \\ S_1^+ \end{pmatrix}, \quad \Phi_{2,3} = \begin{pmatrix} \phi_{2,3}^- \\ -\phi_{2,3}^- \\ S_{2,3} \end{pmatrix}, \quad \Phi_X = \begin{pmatrix} \phi_X^- \\ -\phi_X^- \\ X \end{pmatrix}. \quad (4)$$

While ϕ_1^+ , $\phi_{2,3}^-$ and S_1^+ are electrically charged scalars, the components $\phi_{1,2,3,X}$, $S_{2,3}$ and X are neutral. Therefore, neutral vacuum expectation values (VEVs) are possible in the following directions, $\langle \Phi_1 \rangle^T = (k_1, 0, 0)/\sqrt{2}$, $\langle \Phi_2 \rangle^T = (0, k_2, n)/\sqrt{2}$, $\langle \Phi_3 \rangle^T = (0, k_3, n')/\sqrt{2}$, and $\langle \Phi_X \rangle^T = (0, k_X, n_X)/\sqrt{2}$. However, in order to recover the SM in the low energy limit, we assume the hierarchy $k_{1,2,3,X} \sim v_{SM} \ll n, n', n_X$. Moreover, we consider the particular vacuum structure where $k_2 = n' = k_X = 0$ and $n_X = 0$. The first condition, together with a \mathbf{Z}_2 symmetry (see table I), guarantees the existence of a simple pattern for quark masses, while the second is required to explain the diphoton excess, as will be clear below.² Therefore, the breaking of the gauge groups follows the chain

$$SU(3)_L \otimes U(1)_\mathcal{X} \xrightarrow{n} SU(2)_L \otimes U(1)_Y \xrightarrow{k_{1,3}} U(1)_Q.$$

We note that while the lepton sector does not play an important role for what interests us here, it can nevertheless provide interesting complementary tests to probe the 3-3-1 scale via, for instance, neutrino masses and lepton flavor violation observables [68, 70].

III. PHENOMENOLOGICAL LAGRANGIAN FROM $SU(3)_L \otimes U(1)_\mathcal{X}$

It is instructive to write the Lagrangian in the $SU(2)_L \otimes U(1)_Y$ symmetric phase, i.e., after $SU(3)_L \otimes U(1)_\mathcal{X}$ gets broken at a high-energy scale $\langle S_2 \rangle$. Before

¹ That is, the $SU(3)_L \otimes U(1)_\mathcal{X}$ fundamental representations branch to $SU(2)_L \otimes U(1)_Y$ as: $(\mathbf{3}, \mathcal{X}) \rightarrow (\mathbf{2}, \frac{1}{6} + \mathcal{X}) \oplus (\mathbf{1}, -\frac{1}{3} + \mathcal{X})$ and $(\bar{\mathbf{3}}, \mathcal{X}) \rightarrow (\mathbf{2}, -\frac{1}{6} + \mathcal{X}) \oplus (\mathbf{1}, \frac{1}{3} + \mathcal{X})$.

² Consistency with the minimization conditions of the scalar potential of the model has been checked explicitly.

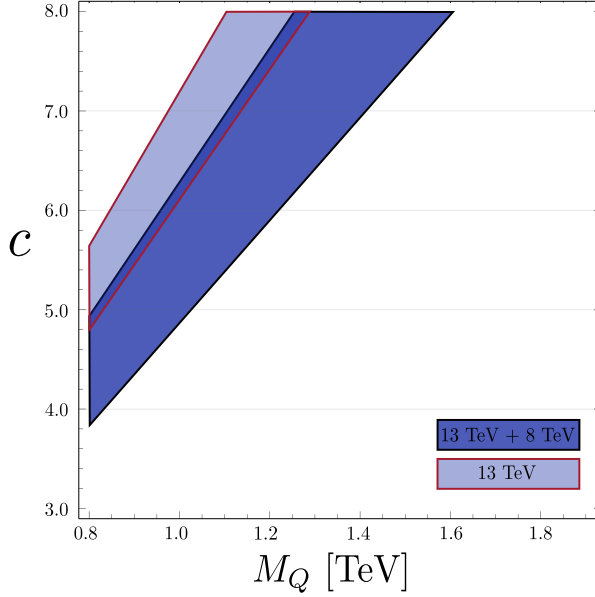


FIG. 1. The coupling between X and the VL quarks, c , as a function of VL quarks mass M_Q . All the points satisfy the bound $M_Q > 800$ GeV and exclusion limits on $\sigma(gg \rightarrow X) \times \text{BR}(X \rightarrow VV)$. The bands correspond to the 95% regions of the combined ATLAS and CMS data for $\sigma(gg \rightarrow X) \times \text{BR}(X \rightarrow \gamma\gamma)$ with 13 TeV $((6.1 \pm 1) \text{ fb [28]})$, dark) and 13 TeV + 8 TeV $((4.4 \pm 1.1) \text{ fb [6]})$, light).

symmetry breaking, the quark Lagrangian reads:

$$\begin{aligned} \mathcal{L}_{\text{quarks}} = & \bar{Q}_L^{1,2} y^u q_R^u \Phi_1^* + \bar{Q}_L^3 \tilde{y}^d q_R^d \Phi_1 + \bar{Q}_L^{1,2} \tilde{y}^d \hat{d}_R \Phi_2^* \\ & + \bar{Q}_L^3 \tilde{y}^u T_R \Phi_2 + \bar{Q}_L^3 \tilde{y}^u q_R^u \Phi_3 + \bar{Q}_L^{1,2} y^d q_R^d \Phi_3^* \\ & + \bar{Q}_L^{1,2} \tilde{y}^d \hat{d}_R \Phi_X + \bar{Q}_L^3 \tilde{y}^u T_R \Phi_X + \text{h.c.}, \end{aligned} \quad (5)$$

where we defined $\hat{d}_R \equiv (D_R, S_R)$, $q_R^u \equiv \{u_R, c_R, t_R\}$ and $q_R^d \equiv \{d_R, s_R, b_R\}$. After $SU(3)_L \otimes U(1)_X$ gets broken to $SU(2)_L \otimes U(1)_Y$, the $SU(3)_L$ triplets split into doublet and singlet representations of $SU(2)_L$. We can write eq. (5) as:

$$\mathcal{L}_{\text{quarks}} = \mathcal{L}_{\text{SM-quarks}} + \mathcal{L}_{S_1^+} + \mathcal{L}_{S_2} + \mathcal{L}_{S_3} + \mathcal{L}_X, \quad (6)$$

where we denote all the terms involving SM quarks and $SU(2)_L$ Higgses (scalar doublets) by $\mathcal{L}_{\text{SM-quarks}}$. This includes for instance terms such as $\bar{q}_L^{1,2} i\sigma_2 H_2^* \hat{d}_R$, with σ_2 the second Pauli matrix and $H_2^\top \equiv (\phi_2 \ \phi_2^-)$. The terms in $\mathcal{L}_{S_1^+}$ involve a charged scalar and so we ignore them here. The terms involving interactions with the neutral singlet S_3 are:

$$\mathcal{L}_{S_3} = c_d'' S_3^* \bar{D}_L q_R^d + c_s'' S_3^* \bar{S}_L q_R^d + c_t'' S_3 \bar{T}_L q_R^u + \text{h.c.}, \quad (7)$$

where we have simplified the notation of the couplings. Since $\langle S_3 \rangle = 0$, the S_3 particle cannot be produced via gluon fusion to quark loops (i.e., there is no mixing be-

tween SM quarks and the new ones³ nor can it decay to photons and thus cannot account for the $X(750)$ resonance). We are then left with two candidates for $X(750)$: X and S_2 , giving the low-energy Lagrangians after $SU(3)_L \otimes U(1)_X$ breaking

$$\mathcal{L}_{S_2} = c_d' S_2^* \bar{D}_L D_R + c_s' S_2^* \bar{S}_L S_R + c_t' S_2 \bar{T}_L T_R + \text{h.c.}, \quad (8)$$

$$\mathcal{L}_X = c_d X^* \bar{D}_L D_R + c_s X^* \bar{S}_L S_R + c_t X \bar{T}_L T_R + \text{h.c.} \quad (9)$$

Again, $c_q^{(\prime)}$ are a simplified notation of the components of the coupling matrices appearing in eq. (5), which we take to be diagonal for simplicity. For instance we defined $\bar{y}_{14}^d \equiv c_d$, $\bar{y}_{25}^d \equiv c_s$, $\bar{y}_{14}^u \equiv c_t$, and $y_{15}^d = y_{24}^d = 0$. From these two possibilities only X is able to reproduce the $X(750)$ resonance as S_2 would imply an unacceptably low $SU(3)_L \otimes U(1)_X$ breaking scale, see sec. IV.

We refer to app. (B) for a detailed discussion of the scalar and gauge boson spectra. It is clear then that \mathcal{L}_X can be seen as an effective Lagrangian (or ‘simplified model’) extending the SM with a neutral iso-singlet scalar, and pairs of vector-like quarks transforming as $(\mathbf{3}, \mathbf{1})_{-1/3}$ for D and S and as $(\mathbf{3}, \mathbf{1})_{2/3}$ for T . X can be produced through gluon fusion and decay to photons via triangle loops involving these new quarks and is therefore our natural candidate for the diphoton resonance. Variants of this effective Lagrangian have been analyzed, see for example [6, 13, 28], and have been shown to be able to account for the diphoton excess. In the next section we derive these results for our specific model.

IV. THE DIPHOTON EXCESS

The X particle can decay to a pair of quarks or gauge bosons. Other decay channels are either kinematically inaccessible or have very suppressed widths, as explained below. Taking all the masses of the new quarks to be M_Q , we can approximately express the widths to gluons

³ We notice that our choice for the vacuum structure of the model guarantees that the exotic quarks that constitute the third components of the $SU(3)_L$ triplets do not mix with the SM states after symmetry breaking. This can be seen in app. (C), where the quark mass matrices are explicitly derived.

and gammas as: [6]

$$\begin{aligned}\Gamma(X \rightarrow gg) &\simeq \frac{2M_X\alpha_s^2}{9\pi} \left(\frac{M_X}{4\pi M_Q} \right)^2 \left| \sum_{i=d,s,t} c_i \right|^2 \\ &= \frac{4M_X\alpha_s^2}{\pi} \left(\frac{M_X}{4\pi n} \right)^2 \left(\frac{c}{c'} \right)^2, \quad (10)\end{aligned}$$

$$\begin{aligned}\Gamma(X \rightarrow \gamma\gamma) &\simeq \frac{M_X\alpha_e^2}{\pi} \left(\frac{M_X}{4\pi M_Q} \right)^2 \left| \sum_{i=d,s,t} c_i Q_i^2 \right|^2 \\ &= \frac{4M_X\alpha_e^2}{3\pi} \left(\frac{M_X}{4\pi n} \right)^2 \left(\frac{c}{c'} \right)^2, \quad (11)\end{aligned}$$

The second equality follows from the assumption of universal couplings $c_i = c$, $c'_i = c'$, and the relation $M_Q = c'n/\sqrt{2}$. It follows that:

$$\begin{aligned}\Gamma(X \rightarrow \gamma\gamma) &= \frac{1}{3} \left(\frac{\alpha_e}{\alpha_s} \right)^2 \Gamma(X \rightarrow gg) \quad (12) \\ &\approx 2.1 \times 10^{-3} \Gamma(X \rightarrow gg),\end{aligned}$$

where we used $\sqrt{4\pi\alpha_s} = 1.07$ at the energy scale $Q = 750$ GeV and $\sqrt{4\pi\alpha_e} = 0.30$ at $Q = 0$. Therefore, in the absence of other decay channels, the branching ratio to photons is $\text{BR}_{\gamma\gamma} \approx 2.1 \times 10^{-3}$. The decay widths into $Z\gamma$ and ZZ via quark loops are always smaller than the $\gamma\gamma$ decay width because they proceed via electroweak mixing; for $SU(2)_L$ -singlet VL fermions the branching ratios are fixed to be: $\text{BR}_{\gamma\gamma} : \text{BR}_{Z\gamma} : \text{BR}_{ZZ} = 1 : 2\tan_W^2 : \tan_W^4$, with $\tan_W \approx 0.55$. These rates are in agreement with ATLAS and CMS bounds. The tree-level decay to ZZ is induced by $Z - Z'$ mixing ($\propto (k/n)^2$), and hence suppressed by the 3-3-1 breaking scale. We found it to be small unless $n \lesssim 1.3$ TeV, which would be in conflict with bounds on Z' direct searches at LHC which are in the multi-TeV range [73, 74]. This also excludes the possibility of a significant contribution of gauge bosons loops to the diphoton signal. Decays to WW are not present because there is no $W - W'$ mixing due to the underlying gauge symmetry.

In fig. (1) we show the variation of the coupling c as a function of the VL quark mass M_Q that satisfies the data. We take the 95% C.L. regions on the combined ATLAS and CMS data using only 13 TeV data from Run II [28] or a combination of 13 TeV and 8 TeV data [6]. On the other hand, in fig. (2) we show for various M_W , the required coupling c' in order to fit the data. The lower-bound on c' translates the bound $M_Q > 800$ GeV on VL quarks. For these figures we have used the exact leading-order relation for $\Gamma(X \rightarrow \gamma\gamma)$ (see app. (A)) instead of the approximation in eq. (10) since $M_Q \sim M_X$. Furthermore, we estimate the production cross-section $\sigma(gg \rightarrow X)$ adapting the results of [6], and we have explicitly checked that these are compatible with those in [4, 28]. We see that the coupling of X with the new

quarks has to be relatively large, ~ 5 for $M_Q \sim 0.8$ TeV, in order to accommodate the 13 TeV data (~ 4 if one considers the combined 13 TeV and 8 TeV data). Still, compared to the case where only one down(up)-type quark is present, the improvement is significant since that would have required couplings as large as $\sim 35(9)$. Also, we note that the physical $X\bar{Q}Q$ vertex is $c/\sqrt{2}$, and not just c , in the perturbative regime.

Another result that can be extracted from fig. (2) is that a hierarchy between c and c' is required in order to explain the diphoton hint. This excludes what could be seen as the minimal possibility of our framework to explain $X(750)$, namely S_2 . Indeed, in a simpler model without the Φ_X triplet, $c = c'$ and the dependence on the coupling is very weak in the decay widths, as can be seen in the approximate relations eq. (10) where it completely disappears. This means that if the signal is interpreted as arising from the decay of a scalar via VL quarks loops, then this scalar cannot be the origin of the VL quarks masses in the context of 3-3-1 models with non-exotic quark charges.

Finally, we comment on the width of the $X(750)$ resonance. Although the ATLAS fit seems to improve if the width is large, $\Gamma \sim 6\%M_X$, current data are perfectly compatible with the existence of a narrow resonance. Nevertheless, we have investigated whether the model could simultaneously account for a large width and the diphoton signal. In principle, the extra decay channel could be provided via the leptonic term $\mathcal{L}_{\text{leptons}} \supset y^s X \bar{N}_L^c N_S$ [70], which is required to generate neutrino masses via the inverse-seesaw mechanism [75].⁴ However, we have found that such term cannot increase the width of the $X(750)$ scalar so as to reach the best-fit value found by ATLAS. Therefore, if future data shows a clear preference for a broad resonance, an extension of our setup will be required.

V. DISCUSSION AND CONCLUSIONS

In this article we have shown that a simple gauge extension of the SM can account for the diphoton excess recently observed by ATLAS and CMS. The gauge structure of the theory requires 3 extra quarks to complete the fundamental $SU(3)_L$ multiplets. These quarks are effectively $SU(2)_L$ singlet VL quarks at low energies. If coupled to a fundamental scalar (anti-)triplet that does not contribute dominantly to their mass (if it takes VEV) then the low-energy iso-singlet explains naturally the $X(750)$ resonance.

The multiplicity of the new quarks, due to the number of families, reduces the severe requirement

⁴ Here N_S denotes a Majorana $SU(3)_L$ singlet (whose majorana mass is the usual ‘ μ term’ of the inverse-seesaw).

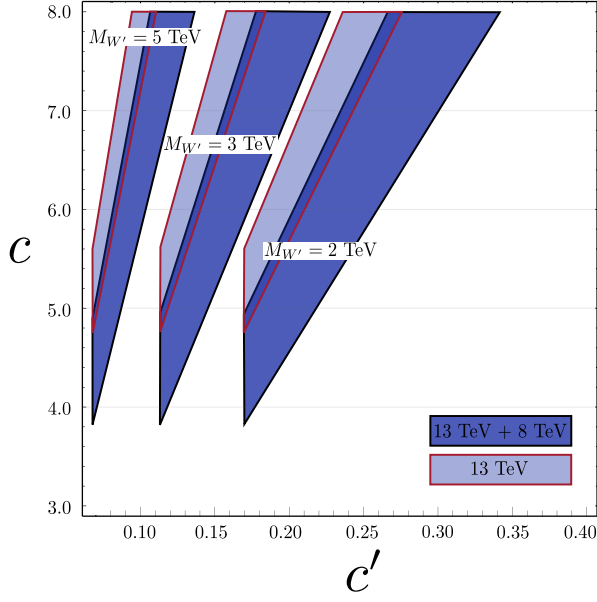


FIG. 2. The coupling between X and the VL quarks, c , as a function of the coupling c' ($M_Q = c'n/\sqrt{2}$) for different masses of the W' gauge boson. All the points satisfy the bound $M_Q > 800$ GeV and exclusion limits on $\sigma(gg \rightarrow X) \times \text{BR}(X \rightarrow VV)$. The bands correspond to the 95% regions of the combined ATLAS and CMS data for $\sigma(gg \rightarrow X) \times \text{BR}(X \rightarrow \gamma\gamma)$ with 13 TeV ((6.1 ± 1) fb [28], dark) and 13 TeV + 8 TeV ((4.4 ± 1.1) fb [6], light).

on the coupling between X and the quarks. This translates as a possible perturbative explanation for the diphoton anomaly. Moreover, we find that X cannot be responsible of the diphoton signal and the breaking of $SU(3)_L \otimes U(1)_\chi$ to the electroweak gauge group at the same time, as that would be excluded by Z' direct searches.

To conclude, we emphasize once again that more data is required in order to fully assess the relevance of the diphoton excess. If confirmed, exciting times will come in the quest of New Physics that expands our current understanding of particles physics. Indeed, X may well be the tip of the iceberg, and future data of LHC will reveal other particles from the UV completion of the theory, possibly in the form of new colored particles and new gauge bosons which are all lying around the TeV scale in our framework.

NOTE ADDED

Shortly after the appearance of our paper, other explanations for the diphoton excess based on the $SU(3)_L \otimes U(1)_\chi$ gauge symmetry were proposed in [76–78]. In contrast to the specific model introduced here, these references consider 3-3-1 models with exotic electric charges, typically leading to slightly lower Yukawa couplings to

explain the diphoton excess.

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Appendix A: $X \rightarrow \gamma\gamma$ width

Following [79, 80], the diphoton decay width of the scalar X to two photons via a loop with particles D , S , T , all with mass $M_Q > M_X/2$, is given by:

$$\Gamma(X \rightarrow \gamma\gamma) = \frac{\alpha_e M_X^3}{512 \pi^3 M_Q^2} \left| N_c \sum_{i=D,S,T} c_i Q_i^2 A_{1/2}(\tau_i) \right|^2, \quad (\text{A1})$$

with $\tau_i = 4M_{Q_i}^2/M_X^2$, $N_c = 3$, Q_i the electric charges of the heavy quarks and

$$A_{1/2}(\tau) = 2\tau^2 (\tau^{-1} + (\tau^{-1} - 1) f(\tau^{-1})) , \quad (\text{A2})$$

$$f(x) = \arcsin^2 \sqrt{x} . \quad (\text{A3})$$

We note that the last expression is valid for a $X(750)$ scalar whose mass is below the kinematic threshold for the production of two heavy Q states.

Appendix B: Scalar potential and mass spectrum

The scalar potential of the model can be written as: ⁵

$$V = \sum_i \mu_i^2 |\Phi_i|^2 + \lambda_i |\Phi_i|^4 + \sum_{i \neq j} \lambda_{ij} |\Phi_i|^2 |\Phi_j|^2 + f(\Phi_1 \Phi_2 \Phi_3 + \text{h.c.}) + \frac{\kappa}{2} [(\Phi_2^\dagger \Phi_X)^2 + \text{h.c.}] , \quad (\text{B1})$$

⁵ For the sake of clarity and simplicity, we consider all the terms involving a single power of $\Phi_{2,X}$ as either absent or with a very small coupling, which is technically natural since this enhances the symmetries of the potential. We also assume CP conservation in the scalar potential.

where $i = 1, 2, 3, X$. The \mathbf{Z}_2 -soft-breaking term, $f\Phi_1\Phi_2\Phi_3$, is required to break accidental symmetries appearing in the scalar potential. Since Φ_2 and Φ_X have identical quantum numbers, all the operators in the scalar potential remain invariant in the exchange $\Phi_2 \leftrightarrow \Phi_X$, however we have assumed in eq. (B1) that operators involving only one power of Φ_X are absent (either because their couplings are so small that they can be ignored, or because of a symmetry which distinguishes Φ_2 from Φ_X , e.g. a parity) in order to avoid tree-level mixing between X and the SM Higgses and allow for the particular vacuum solution with $n_X = 0$. Assuming $k_1 \sim k_3 \equiv k \ll n$ and universal couplings, we find that the mass spectrum of the charged scalars is given at leading order as: ⁶

$$M^2(\phi_2^\pm) = 0, \quad (\text{B2})$$

$$M^2((\phi_1^\pm + \phi_3^\pm)/\sqrt{2}) = 0, \quad (\text{B3})$$

$$M^2((\phi_1^\pm - \phi_3^\pm)/\sqrt{2}) \sim \frac{1}{\sqrt{2}} f n, \quad (\text{B4})$$

$$M^2(S_1^\pm) \sim \sqrt{2} f n, \quad (\text{B5})$$

$$M^2(\phi_X^\pm) = \mu_X^2 + \lambda k^2. \quad (\text{B6})$$

The masses of the neutral CP-even scalars are, up to corrections of $\mathcal{O}(k^2)$:

$$M^2(\Re(\phi_1 + \phi_3)/\sqrt{2}) \sim (2\lambda + \sqrt{2} \frac{f}{n} - \frac{1}{2} \frac{f^2}{\lambda n^2}) k^2, \quad (\text{B7})$$

$$M^2(\Re(\phi_1 - \phi_3)/\sqrt{2}) \sim \frac{1}{\sqrt{2}} f n, \quad (\text{B8})$$

$$M^2(\Re\phi_2) = 0, \quad (\text{B9})$$

$$M^2(\Re S_2) \sim 2\lambda n^2, \quad (\text{B10})$$

$$M^2(\Re S_3) \sim \sqrt{2} f n, \quad (\text{B11})$$

$$M^2(\Re\phi_X) = \mu_X^2 + \lambda k^2. \quad (\text{B12})$$

On the other hand, the masses of the neutral CP-odd scalars at leading order are given as:

$$M^2(\Im(\phi_1 + \phi_3)/\sqrt{2}) \sim \sqrt{2} f n, \quad (\text{B13})$$

$$M^2(\Im(\phi_1 - \phi_3)/\sqrt{2}) = 0, \quad (\text{B14})$$

$$M^2(\Im\phi_2) = 0, \quad (\text{B15})$$

$$M^2(\Im S_2) = 0, \quad (\text{B16})$$

$$M^2(\Im S_3) \sim \frac{1}{\sqrt{2}} f n, \quad (\text{B17})$$

$$M^2(\Im\phi_X) = \mu_X^2 + \lambda k^2. \quad (\text{B18})$$

Finally, the state $\Re X \equiv X$ has a mass:

$$\begin{aligned} M^2(\Re X) &= M_X^2 \\ &= \mu_X^2 + \frac{1}{2} (\lambda_{1X} + \lambda_{3X}) k^2 + \frac{1}{2} \kappa n^2 \end{aligned} \quad (\text{B19})$$

and does not mix with the other CP-even scalars. The mass of its CP-odd counterpart, $\Im X$, can be independently set with a proper choice of the κ quartic coupling, as one finds $M^2(\Re X) - M^2(\Im X) = \kappa n^2$. The massless scalars found in the above equations correspond to the degrees of freedom ‘eaten-up’ by the charged and neutral gauge bosons, respectively, which acquire the following masses:

$$m_W^2 = \frac{1}{2} g_2^2 k^2, \quad (\text{B20})$$

$$m_{W'}^2 = \frac{1}{4} g_2^2 n^2, \quad (\text{B21})$$

$$m_Z^2 = \frac{g_2^2(4g_1^2 + 3g_2^2)}{2(g_1^2 + 3g_2^2)} k^2, \quad (\text{B22})$$

$$m_{Z'}^2 = \frac{1}{9} (g_1^2 + 3g_2^2) n^2, \quad (\text{B23})$$

$$m_{X_G}^2 = m_{Y_G}^2 = \frac{1}{4} g_2^2 n^2. \quad (\text{B24})$$

with $g_2(g)$ being the coupling constant of $SU(3)_L$ ($U(1)_X$). Notice that since S_2 is singlet under the $SU(2)_L$ subgroup contained in $SU(3)_L$, the VEV n will control the four new gauge bosons masses and break $SU(3)_L$ to $SU(2)_L$. On the other hand, $SU(2)_L \otimes U(1)_Y$ is broken at the electroweak scale by the k_1 and k_3 VEVs down to the electromagnetic $U(1)_Q$ symmetry. For $f \sim n$ all the scalars of the model are naturally heavy, except one state that we can identify with the SM Higgs boson, i.e., $H \equiv (\phi_1 + \phi_3)/\sqrt{2}$, in good approximation. Indeed, its couplings to the fermions confirm that the state H is the one that gives mass to the SM fermions.

Appendix C: Quark masses

The quark Lagrangian in eq. (5) leads to the following mass matrices:

$$M_d = -\frac{1}{\sqrt{2}} \begin{pmatrix} y_{11}^d k_2 & y_{12}^d k_2 & y_{13}^d k_2 & 0 & 0 \\ y_{21}^d k_2 & y_{22}^d k_2 & y_{23}^d k_2 & 0 & 0 \\ \tilde{y}_{11}^d k_1 & \tilde{y}_{12}^d k_1 & \tilde{y}_{13}^d k_1 & 0 & 0 \\ 0 & 0 & 0 & \bar{y}_{14}^d n & \bar{y}_{15}^d n \\ 0 & 0 & 0 & \bar{y}_{24}^d n & \bar{y}_{25}^d n \end{pmatrix}, \quad (\text{C1})$$

$$M_u = -\frac{1}{\sqrt{2}} \begin{pmatrix} y_{11}^u k_1 & y_{12}^u k_1 & y_{13}^u k_1 & 0 \\ y_{21}^u k_1 & y_{22}^u k_1 & y_{23}^u k_1 & 0 \\ \tilde{y}_{11}^u k_2 & \tilde{y}_{12}^u k_2 & \tilde{y}_{13}^u k_2 & 0 \\ 0 & 0 & 0 & \bar{y}_{14}^u n \end{pmatrix}. \quad (\text{C2})$$

The \mathbf{Z}_2 symmetry, see table I, is introduced so that the SM quarks and the new ones are independent of each other and can be adjusted individually to easily obtain a realistic quark sector and heavy exotic quarks at the same time.

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